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Scheme Independence at First Order Phase Transitions and the Renormalisation Group

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Abstract

We analyse approximate solutions to an exact renormalisation group equation with particular emphasis on their dependence on the regularisation scheme, which is kept arbitrary. Physical quantities related to the coarse-grained potential of scalar QED display universal behaviour for strongly first-order phase transitions. Only subleading corrections depend on the regularisation scheme and are suppressed by a sufficiently large UV scale. We calculate the relevant coarse-graining scale and give a condition for the applicability of Langer's theory of bubble nucleation.

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1. A promising tool for the investigation of non-perturbative effects in quantum field theory is the Exact Renormalisation Group (ERG) [1]-[9]. Conceptually, ERG equations are quite appealing as they relate (effective) degrees of freedom at different length scales. *Solving* an ERG equation is much more intriguing. As it couples an infinite number of operators finding any exact solution would seem to be impossible. Thus, any practical computation has in addition to face the problem of finding an appropriate truncation/approximation of the infinite dimensional space of operators. However, very little is known about the regularisation scheme (RS) dependence of approximate solutions. Clearly, an approximation scheme has to be discarded in case that a change of the RS either changes the *qualitative* behaviour of the solution, or introduces *large quantitative* corrections.[1] Therefore, it seems very useful to investigate an example where the RS dependence can be made explicit. In this Letter we study the RS dependence of physical quantities related to first order phase transitions at the example of scalar quantum electrodynamics in Euclidean space time [10, 11]. We follow the approach advocated in [2, 3]. Our conclusions agree with recent numerical investigations [12, 13, 14]. In addition, they provide further evidence for the viability of the approximations and yield a condition for Langer's theory of bubble nucleation [15] to be applicable.

2. The ERG equation [2] reads (with $t = \ln k$) for bosonic fields Φ

$$\frac{\partial}{\partial t}\Gamma_k[\Phi] = \frac{1}{2}\mathrm{Tr}\left\{\left(\Gamma_k^{(2)}[\Phi] + R_k\right)^{-1}\frac{\partial R_k}{\partial t}\right\} \tag{1}$$

where the length scale k^{-1} can be interpreted as a coarse-graining scale [2, 3]. Eq. (1) relates the microscopic action $S[\Phi] = \lim_{k\to\infty}\Gamma_k[\Phi]$ with the corresponding macroscopic (effective) action $\Gamma[\Phi] = \lim_{k\to 0}\Gamma_k[\Phi]$, the generating functional of 1PI Green functions. The right hand side of eq. (1) involves the regulator function R_k and the second functional derivative of the effective average action w.r.t. the fields. The trace stands for summation over all indices and integration over all momenta. In momentum space, R_k is a function of momentum squared q^2, and we will write it, for dimensional reasons, as

$$R_k(q^2) = q^2 g\left(\frac{q^2}{k^2}\right) \ . \tag{2}$$

In order to regularise zero modes of the second functional derivative of Γ_k, the following conditions on the RS have to be met:

(i) $\lim_{q^2\to 0} R_k(q^2) > 0.$

(ii) $\lim_{k\to 0} R_k(q^2) = 0.$

(iii) $\lim_{k\to\infty} R_k(q^2) \to \infty.$

[1] A recent letter [9] showed that even critical exponents can depend rather strongly on the RS.

Condition (i) ensures that R_k acts like an additional (possibly momentum dependent) mass term for small momenta and hinders infrared divergencies in the case of massless modes. (In the sharp cutoff limit [1] it reads $\lim_{q^2\to 0} R_k(q^2) = \infty$.) Condition (ii) allows to recover the usual effective action in the limit $k \to 0$ as any dependence on R_k drops out. Condition (iii) ensures that the classical action is obtained for $k \to \infty$. For any practical applications, condition (iii) is weakened, replacing $k \to \infty$ by $k \to \Lambda$ (where Λ is some UV cutoff scale), and both Λ and $R_\Lambda(q^2)$ have to be much larger than any other physical scale in the theory (typically $R_\Lambda \sim \Lambda^2$). A general class of smooth RSs is given for arbitrary $b \geq 1$ by [2]

$$g(y) = (\exp y^b - 1)^{-1} \ . \tag{3}$$

Sometimes it is also very convenient to use a simple power-like regulator [7, 12]

$$g(y) = 1/y^b \ . \tag{4}$$

It can be shown that the limit $b \to \infty$ of eq. (3) and eq. (4) corresponds to the sharp cutoff limit [1].

3. We now specify the flow equation for the Abelian Higgs model in d dimensions (refering the reader to [3] for their derivation), and detail the approximations used. Our Ansatz for the relevant part of Γ_k uses the first term(s) of a derivative expansion

$$\Gamma_k[\varphi, A] = \int d^d x \left\{ U_k(\bar\rho) + \frac{1}{4} F_{\mu\nu}F_{\mu\nu} + (D_\mu[A]\varphi)^* D_\mu[A]\varphi \right\} \tag{5}$$

and neglects all higher derivative terms. Here, $F_{\mu\nu}$ denotes the usual Abelian field strength, $\bar\rho = \varphi^*\varphi$, $D_\mu = \partial_\mu + i\bar e A_\mu$ the covariant derivative and $\bar e$ the (dimensionful) Abelian charge. We neglect the anomalous scaling of the gauge- and scalar kinetic terms and the running of the Abelian charge, which is a valid approximation near the Gaussian fixed point where the anomalous dimensions are known to be small [13, 14]. It follows the flow equation for the coarse-grained potential U_k as

$$\frac{dU_k(\bar\rho)}{2v_d k^{d-1} dk} = \ell_0^d\left(\frac{U_k'(\bar\rho) + 2\bar\rho\, U_k''(\bar\rho)}{k^2}\right) + \ell_0^d\left(\frac{U_k'(\bar\rho)}{k^2}\right) + (d-1)\ell_0^d\left(\frac{2\bar e^2\bar\rho}{k^2}\right) , \tag{6}$$

where $v_d^{-1} = 2^{d+1}\pi^{d/2}\Gamma[\frac{d}{2}]$ and the primes denote the derivatives w.r.t. $\bar\rho$. We readily identify the different terms on the r.h.s. as the contributions from the massive scalar/massless scalar/gauge field fluctuations respectively. The threshold functions

$$\ell_n^d(w) = -\int_0^\infty dy\, y^{\frac{d}{2}+1} \frac{(n+\delta_{0,n})\, g'(y)}{[(1+g(y))y + w]^{n+1}} \tag{7}$$

encode the dependence on the RS. The integration over $y = q^2/k^2$ (momentum squared in units of k) stems from the Tr on the r.h.s. of (1). Figure 1 displays the function $\ell_2^4(\omega)$ for different RS. Note that these functions already differ in their order of magnitude for

moderate values of ω. As long as the quartic scalar self coupling $\bar{\lambda}$ is small compared to the gauge coupling $\bar{e}^2$ we will in addition neglect the fluctuations from the scalar field. It is known that this is the case for the Coleman-Weinberg phase transition in four dimensions [10, 13] and for a part of the phase diagram of normal superconductors in three dimensions where the phase transition is first order [11, 12, 14]. The flow equation then reads

$$\frac{dU_k(\bar{\rho})}{dk} = 2(d-1)v_d k^{d-1}\ell_0^d\left(\frac{2\bar{e}^2\bar{\rho}}{k^2}\right) . \tag{8}$$

This approximation can be controlled comparing the gauge-field-induced contributions to the scale dependence of U_k'' with those that have been neglected. Note also that eq. (8) still allows for arbitrarily high (scalar) couplings. No assumptions on the functional form of the potential are made, and we do in particular not assume $U_k(\bar{\rho})$ to be a local polynomial in $\bar{\rho}$ (which it is not). However, eq. (8) will not allow an investigation of the flattening of the inner part of the potential in the limit $k \to 0$, which is known [16] to be triggered by the terms $\sim \ell_0^d(U'/k^2) + \ell_0^d(U'/k^2 + 2\bar{\rho}U''/k^2)$. It follows, that the approximation becomes invalid in the non-convex part of the potential at some *flattening scale* $k_\flat > 0$. We will come back to this point later. For the time being we will study solutions to eq. (8) and their RS dependence with an initial potential at some UV scale Λ given by $U_\Lambda(\bar{\rho}) = m_\Lambda^2\bar{\rho} + \frac{1}{2}\lambda_\Lambda\bar{\rho}^2$. The coarse grained potential obtains as $U_{k,\Lambda}(\bar{\rho}) = U_\Lambda(\bar{\rho}) + \Delta(\bar{\rho})$, where

$$\Delta(\bar{\rho}) = \int_k^\Lambda d\bar{k}\int_0^\infty dy\, \frac{2(d-1)v_d y^{\frac{d}{2}+1}g'(y)\,\bar{k}^5}{[g(y)+1]y\;\bar{k}^2 + 2\bar{e}^2\bar{\rho}} . \tag{9}$$

stems from integrating out fluctuations between Λ and k. Three different mass scales are related to the coarse grained potential: The mass of the scalar field in the regime with spontaneous symmetry breaking (SSB), m, the scalar mass in the symmetric regime, m_s, and the mass of the gauge field in the SSB regime, M. They are given by

$$\begin{aligned} m_s^2(k) &= U_k'(\bar{\rho}=0) \\ m^2(k) &= 2U_k''(\bar{\rho}_0)\bar{\rho}_0 \\ M^2(k) &= 2\bar{e}^2\bar{\rho}_0 \ , \end{aligned} \tag{10}$$

where $\bar{\rho}_0$ denotes the location of the minimum in the SSB regime. In the limit $k \to 0$ these masses correspond to the two-point function at vanishing external momentum.

4. To be explicit, we switch to four dimensions, although the following discussions can be made for any d.[2] We will focus in the sequel on a) a mass relation, b) the degenerate (critical) potential, and c) the surface tension.

a) We start with the scale dependent *mass relation*

$$\frac{m^2}{M^2} + 2\frac{m_s^2}{M^2} = F\left(\frac{M}{\Lambda}, \frac{k}{M}, e^2\right) \tag{11}$$

[2]As the gauge coupling in $d=4$ is dimensionless, we will substitute $\bar{e}^2 \to e^2$.

with F vanishing at $k = \Lambda$.[3] F is related to $\Delta(\bar\rho)$ by

$$F = \frac{1}{e^2}\left\{\Delta''(\bar\rho_0) + \frac{\Delta'(0) - \Delta'(\bar\rho_0)}{\bar\rho_0}\right\} \ .$$

Using eq. (9) and performing the $\bar k$-integration first gives

$$F = \frac{3e^2}{8\pi^2}\,\mathcal{I}_g\left[\frac{2 + 3p^2s^2}{2[1 + p^2s^2]^2} - \left(s \leftrightarrow \frac{1}{r}\right)\right] \ , \tag{12}$$

where we have used $r = M/\Lambda$ and $s = k/M$. $p^2 = y + yg(y)$ denotes the inverse effective propagator in units of k^2. We have also introduced the linear integral operator

$$\mathcal{I}_g[f] = -2\int_0^\infty dy\,\frac{g'(y)\,f(y)}{[1 + g(y)]^3}\,. \tag{13}$$

that provides a RS dependent measure in momentum space. As the argument of $\mathcal{I}_g$ in eq. (12) depends on momenta, it is understood that the mass relation explicitely depends on the RS. Note, however, that $\mathcal{I}_g[1] = 1$ is *independent* of g, as long as g fullfills the conditions (i) and (ii). Furthermore, in the limit $k \to 0$ and $\Lambda \to \infty$ we obtain

$$\frac{m^2}{M^2} + 2\frac{m_s^2}{M^2} = \frac{3e^2}{8\pi^2} \tag{14}$$

which is independent of both the RS and the initial conditions.[4] Of course, for general k, Λ the r.h.s. of eq. (12) is not universal and will depend on the precise shape of $g(y)$. We have displayed the mass relation for different RS and $\Lambda \to \infty$ in figure 2, with H given by $3e^2H = 8\pi^2F$. H describes the cross-over from a "classical" ($H \approx 0$) region, where fluctuations are of no importance, to a "coarse-grained" ($H \approx 1$) region, where fluctuations have already been integrated out. The precise form of the cross-over around $k \approx M$ depends on the RS provided. Corrections due to a finite Λ introduce RS dependent terms. Expanding eq. (12) for $k = 0$ in powers of $(M/\Lambda)^2$, we obtain

$$F = \frac{3e^2}{8\pi^2}\sum_{n=0}^{\infty}(-)^n\left(1 + \frac{n}{2}\right)a_n^g\ r^{2n}\,. \tag{15}$$

The leading correction is already quadratically suppressed. The coefficients a_n^g are the (RS dependent) n^{th} moments of $1/p^2$ w.r.t. the measure $\mathcal{I}_g$,

$$a_n^g = \mathcal{I}_g\left[p^{-2n}\right]\,. \tag{16}$$

With g as in eq. (4) they read $a_n^g = 2\Gamma[1 + \frac{n}{b}]\Gamma[2 + n - \frac{n}{b}]/\Gamma[3 + n]$. For large n, they decay at least with $1/n$. A similar behaviour is observed numerically in the case of the

[3] Any potential quadratic in $\bar\rho$ with $M \neq 0$ obeys $m^2 + 2m_s^2 = 0$. This relation is best suited to study properties of the coarse-grained potential that are *independent* of the initial conditions.

[4] An analogous result was found in [11].

exponential regulator eq. (3). The first subleading coefficient a_1^g ranges only between 1/3 and 2/3 for g from both eq. (3) and (4) and with $1 \leq b \leq \infty$. It is quite remarkable that these coefficients are rather insensitive against a change of the RS. This is related to the fact that only an *average* of the RS enters in a_n^g: Assuming that $g(y)$ is monotonous, we can rewrite eq. (16) as a convolution over g,

$$a_n^g = 2\int_0^\infty dg \frac{1}{(1+g)^{3+n}}\ \frac{1}{y(g)^n}\ .$$

The first function in the integrand vanishes for large g and has its maximum for small g, whereas the second function vanishes for small g [condition (ii)] and grows large for large g [condition (iii)]. Thus, the integrand is peaked and the integral will get contributions over a rather broad range of values for g, which explains partly the weak RS dependence of the a_n^g and is in contrast to the RS dependence encountered in [9].
The *coarse-graining scale* $k_\natural$ can be estimated from eq. (12), that, expanded in $s = k/M$, gives $H = 1 - \frac{1}{2}a_{-1}^g s^2 + \mathcal{O}(s^4)$. For g from eq. (3), the coefficient $a_{-1}^g = 2\Gamma[1+\frac{1}{b}]$ ranges between 1.77 and 2 with $1 \leq b \leq \infty$. The plateau is reached at the 1% level as soon as the coarse-graining scale is about 10% of the gauge-field mass,

$$k_\natural \approx 0.1\, M\ . \tag{17}$$

b) The *critical potential* $U_{k,\Lambda}^{\rm crit}$ corresponds to the one with degenerate minima, $U_{k,\Lambda}^{\rm crit}(0) = U_{k,\Lambda}^{\rm crit}(\bar\rho_0)$. The initial conditions, that lead to a degenerate potential with v.e.v. $\bar\rho_0$ at some scale k are uniquely specified as

$$\begin{aligned} m_{\Lambda,\rm c}^2 &= \Delta'(\bar\rho_0) - \frac{2}{\bar\rho_0}\left\{\Delta(\bar\rho_0) - \Delta(0)\right\} \\ \lambda_{\Lambda,\rm c} &= -\frac{2}{\bar\rho_0^2}\left\{\bar\rho_0\Delta'(\bar\rho_0) - \Delta(\bar\rho_0) + \Delta(0)\right\}\ , \end{aligned} \tag{18}$$

and we obtain

$$U_{k,\Lambda}^{\rm crit}(\bar\rho) = \frac{\bar\rho}{\bar\rho_0}[\Delta(\bar\rho_0) - \Delta(0)]\left(\frac{\bar\rho}{\bar\rho_0} - 2\right) + \Delta'(\bar\rho_0)\bar\rho\left(1 - \frac{\bar\rho}{\bar\rho_0}\right) + \Delta(\bar\rho)\ . \tag{19}$$

Using eq. (9) and eq. (13), the critical potential follows (apart from an irrelevant field-independent constant) as

$$U_{k,\Lambda}^{\rm crit}(\bar\rho) = 6v_4e^4\left\{\mathcal{I}_g\left[\bar\rho^2\,\ln\frac{2e^2\bar\rho + p^2k^2}{M^2+p^2k^2} + \frac{M^2\bar\rho(\bar\rho_0 - \bar\rho)}{M^2+p^2k^2}\right] - (k\leftrightarrow\Lambda)\right\}. \tag{20}$$

Again in the limit $k \to 0, \Lambda \to \infty$ any RS dependence of the critical potential drops out, and, using $z = \bar\rho/\bar\rho_0$, we are left with the potential

$$U_{0,\infty}^{\rm crit}(\bar\rho) = \frac{3M^4}{64\pi^2}\left\{z^2\ln z + z\,(1-z)\right\} \tag{21}$$

and $m^2/M^2 = 6v_4 e^2$. Corrections due to a finite Λ can be expanded as a power series in $(M/\Lambda)^2$. For $k=0$ we obtain

$$U_{0,\Lambda}^{\rm crit} = U_{0,\infty}^{\rm crit} - \frac{3}{2} v_4 M^4 \sum_{n=0}^{\infty} a_{n+1}^g \; p_n(z) z(1-z)^2 \frac{(-)^n r^{2n+2}}{n+\delta_{n,0}} \tag{22}$$

where p_n are n^{th}-order polynomials in $\bar\rho/\bar\rho_0$,

$$p_n(z) = \sum_{m=0}^{n} (m+1) z^{n-m} \; .$$

c) The *surface tension* σ, defined as

$$\sigma_{k,\Lambda}(\bar\rho_0) = \int_0^{\bar\varphi_0} d\varphi \sqrt{2U_{k,\Lambda}^{\rm crit}(\bar\rho)} \tag{23}$$

is sensitive to the shape of the critical potential. We expect on dimensional grounds that it scales like $\sigma \sim \bar\rho_0^{\ 3/2}$, and with eq. (21) we obtain

$$\sigma = \sqrt{3v_4}\, a_0 e^2 \, \bar\rho_0^{\ 3/2} = \frac{3\, a_0}{32\pi^2} \frac{M^4}{m} \; . \tag{24}$$

The coefficient $a_0 = \int_0^1 dx \sqrt{1-x+x\ln x} \approx 0.42$ encodes the shape of the critical potential. Defining $\delta\sigma = \sigma - \sigma_{0,\Lambda}$ we again find that RS dependent terms are suppressed by powers in $(M/\Lambda)^2$,

$$\frac{\delta\sigma}{\sigma} = \sum_{n=1}^{\infty} (-)^{n+1} \; a_n^g \, \frac{a_n}{a_0} \, r^{2n} \; . \tag{25}$$

The expansion coefficients depend on the shape of the critical potential encoded in the numerical factors a_n with $a_1 = \int_0^1 dx (1-x)^2/\sqrt{x \; \ln x + 1 - x} \approx 0.30$ (and similar expressions for the higher terms $a_2 \approx 0.24, a_3 \approx 0.18, \ldots$), and on the scheme dependent numbers a_n^g. In figure 3 we have displayed both the v.e.v. $\bar\rho_0$ and the surface tension as functions of m/M.

5. The small-momentum fluctuations of the scalar field become important for the non-convex part of the potential. From eq. (6) we can estimate this scale through $k_\flat^2 \approx \max\{-U', -U' - 2\bar\rho U''\}$ and obtain with eq. (21)

$$k_\flat^2 \approx 3 v_4 e^2 M^2 \; . \tag{26}$$

Evaluating the mass relation eq. (12) at $k = k_\flat$ instead of $k=0$ induces subleading corrections $\sim v_4 e^2$ to eq. (14). This is consistent with an estimation using the partial differential equation (6). Defining $\rho = k^{2-d}\bar\rho$, $u(\rho) = k^{-d} U(\bar\rho)$, $e^2 = k^{d-4}\bar e^2$ and expanding the r.h.s. of the flow equation for $u(\rho)$ up to linear order in $u'(\rho)$, we obtain

$$\frac{\partial u}{\partial t} = -d\, u + \Big[(d-2)\rho - 4 v_d \ell_1^d(0)\Big] \frac{\partial u}{\partial \rho} + 4 v_d \ell_0^d(0) + 2(d-1) v_d \ell_0^d(2e^2\rho) \, . \tag{27}$$

The corresponding mass relation eq. (12) obtains a more complicated e^2-dependence that reads in the sharp cutoff limit $m^2/M^2 + 2m_s^2/M^2 = 3e^2/8\pi^2/(1+2v_4e^2)^2$ [12, 18].

It remains to be shown that the surface tension has already reached a plateau at the coarse-graining scale $k_\natural > k_\flat$ before the scalar fluctuations are relevant. This is an important prerequisite for Langer's theory of bubble nucleation [15] to be applicable, and of relevance for the electroweak phase transition [19, 20]. In figure 4 the surface tension is calculated as a function of k/M for fixed v.e.v. $\bar\rho_0$ and different RSs. The initial conditions are chosen as to give a degenerate potential with v.e.v. $\bar\rho_0$ at some scale k with the corresponding surface tension σ_k. In the limit $k \to 0$ the critical initial values eq. (18) are independent of k. Decreasing k shows that σ initially increases until it reaches a constant value. The ratio of flattening to coarse-graining scale obtains from eq. (17) and eq. (26) as

$$k_\flat^2/k_\natural^2 \approx e^2 \tag{28}$$

and is effectively RS independent.[5] The condition for a coarse-grained surface tension to be a well-defined quantity ($k_\flat/k_\natural < 1$) is automatically fullfilled with $e^2 \ll 1$. Ultimately, this is related to the strength of the phase transition, being strongly first order. For weakly first order transitions we have to expect that $k_\flat/k_\natural \sim \mathcal{O}(1)$ or even larger, and defining a coarse-grained surface tension becomes ambiguous.[6]

Our results are by no means specific to the four dimensional case. The main difference in three dimensions comes from the fact that the gauge coupling has the dimension of a mass. Furthermore, the scale dependence of $\bar e^2$ is in general no longer negligible and a non-trivial fixed point structure governs the flow equation [12, 14, 17]. Our approximations remain valid as long as the gauge coupling stays in the vicinity of the Gaussian fixed point. The mass relation eq. (14) then becomes

$$m^2 + 2m_s^2 = \frac{1}{2\pi}\bar e^2 M \tag{29}$$

and the critical potential eq. (21) reads

$$U = \frac{M^3}{12\pi}\left[z(1+z) - 2z^{3/2}\right] \ . \tag{30}$$

Note the appearence of the non-analytic term $\sim |\varphi|^3$ in the potential. The condition eq. (28) generalises to

$$k_\flat^2/k_\natural^2 \approx \frac{\bar e^2}{M} \ . \tag{31}$$

With $\bar e^2/M$ being the perturbative expansion parameter it follows that Langer's theory is viable within the perturbative regime. A full discussion where the running gauge coupling is also taken into account will be given elsewhere [22].

[5] The uncertainty in defining the scales $k_\flat$ and $k_\natural$ is much larger than the RS dependence.

[6] See, however, the discussion for 3d matrix models in [21], where also a result analogous to our eq. (31) is given.

6. In summary, we calculated physical quantities at a photon induced first order phase transition and studied their RS (in-)dependence. The main RS dependence enters through the finite UV scale $\Lambda < \infty$ and the scale $k_\flat > 0$, both of which introduce terms proportional to powers of $k_\flat/M$ or M/Λ, with M being a typical mass scale of the theory. The related expansion coefficients a_n^g show a very weak RS dependence which indicates the stability of the approximation used. Especially, as the RS dependent terms introduce only marginal quantitative corrections, none of the physical conclusions are affected by them. This is of immediate relevance for the sophisticated numerical investigations presented in [20, 21]. Note also that these RS dependent terms would not show up at a second order phase transition, where a scaling (*i.e.* k-independent) solution is obtained. It emerged that the concept of a coarse-grained surface tension based on the seperation of low- and high energy modes is only viable for sufficiently strongly first order phase transitions. The criterion for the validity of the standard treatment of bubble nucleation is RS independent.

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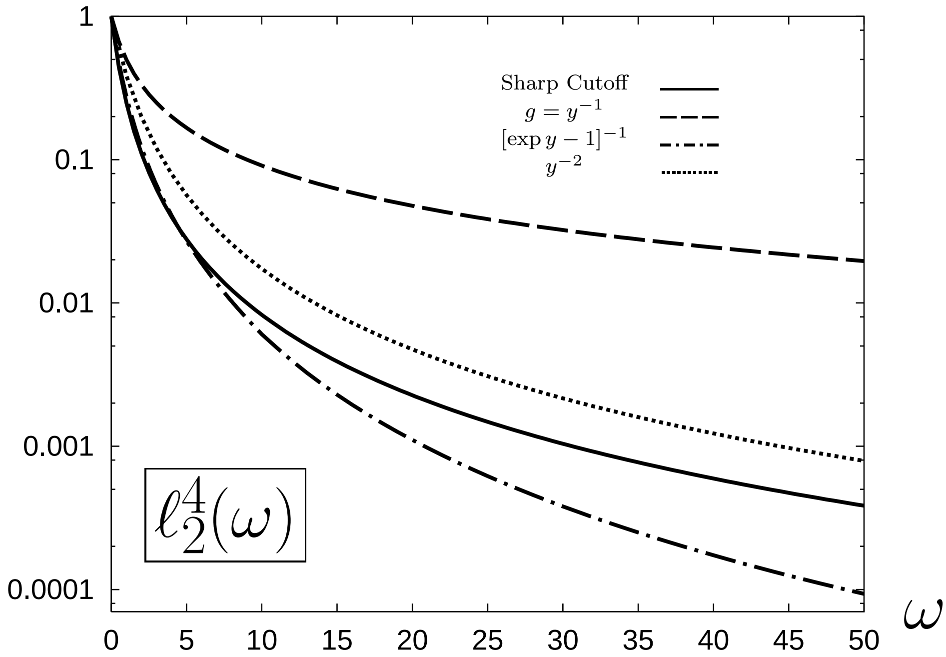

Figure 1: The threshold function $\ell_2^4(\omega)$ for different RS.

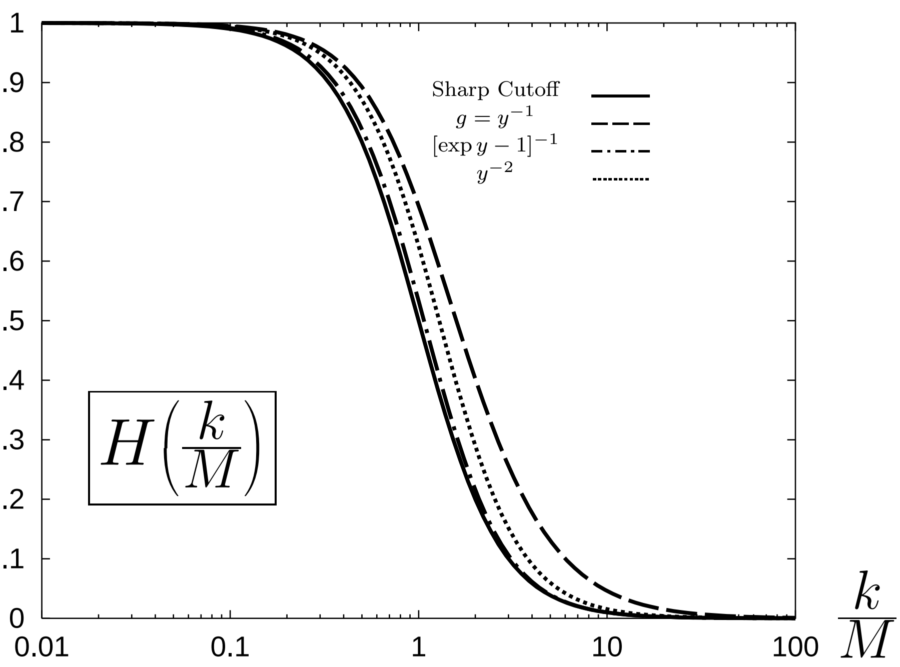

Figure 2: The mass relation for different RS.

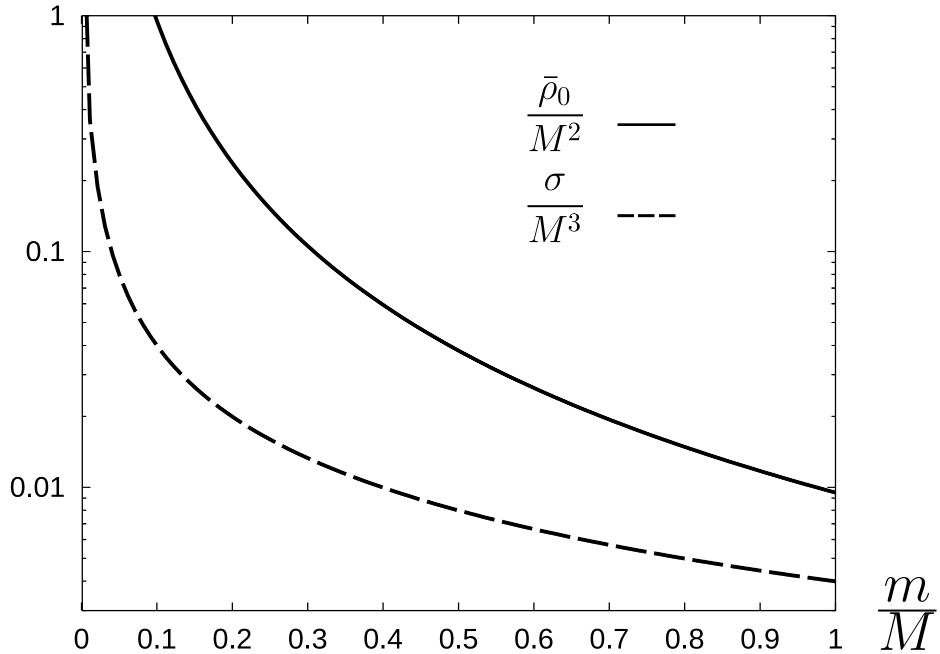

Figure 3: The surface tension σ and the v.e.v. $\bar{\rho}_0$ as functions of m/M.

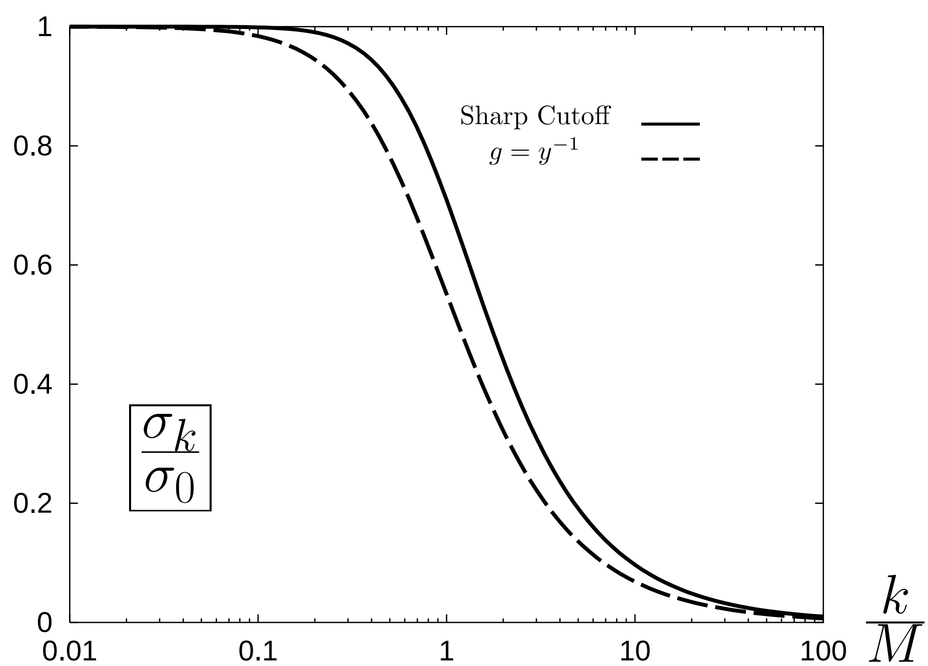

Figure 4: The surface tension for fixed v.e.v. $\bar{\rho}_0$ as a function of k/M for different RS in the limit $\Lambda \to \infty$.